

# Swift J1644+57: A White Dwarf Tidally Disrupted by a $10^4 M_{\odot}$ Black Hole?

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## ABSTRACT

We propose that the remarkable object Swift J1644+57, in which multiple recurring hard X-ray flares were seen over a span of several days, is a system in which a white dwarf was tidally disrupted by an intermediate mass black hole. Disruption of a white dwarf rather than a main sequence star offers a number of advantages in understanding the multiple, and short, timescales seen in the light curve of this system. In particular, the short internal dynamical timescale of a white dwarf offers a more natural way of understanding the short rise times ( $\sim 100$  s) observed. The relatively long intervals between flares ( $\sim 5 \times 10^4$  s) may also be readily understood as the period between successive pericenter passages of the remnant white dwarf. In addition, the expected jet power is larger when a white dwarf is disrupted. If this model is correct, the black hole responsible must have mass  $\lesssim 10^5 M_\odot$ .

*Subject headings:* accretion, black holes, white dwarfs

## 1. Introduction

On 28 March 2011, the *Swift* Burst Alert Telescope (BAT) detected a most unusual object, Swift J164449.3+573451 (Levan *et al.* 2011; Burrows *et al.* 2011). Although in many ways this object (whose name we abbreviate to Swift J1644+57) initially appeared to resemble a classical  $\gamma$ -ray burst, its light curve soon showed that it was quite different. Still bright more than  $6 \times 10^6$  s after the initial trigger (see Fig. 1), in its initial activity it exhibited repeated extremely short timescale flares (see Fig. 2). After holding roughly steady for  $\simeq 700$  s, the flare causing the BAT trigger rose a factor of 10 in flux over the next  $\simeq 400$  s. Less than 1000 s later, the flux had fallen by a factor of 20, and  $\sim 10^4$  s after that, although still detectable, the flux was only 0.5% of what it had been at the peak. Most surprisingly, there was a comparable flare  $\simeq 50000$  s later, similarly lasting for only  $\sim 1000$  s, and a third, slightly brighter than the first,  $\simeq 60000$  s after that. Like the first two, the third flare continued to show very short timescale variation.

Several more brief flares of comparable brightness followed, likewise separated by periods of flux two orders of magnitude weaker. Gradually the duration of the flares stretched and their amplitudes diminished, until after  $\simeq 2 \times 10^5$  s the system began a long, gradual decline in flux in which it has remained bright enough to be detected for more than  $1.5 \times 10^7$  s. Thus, there appear to be a number of characteristic timescales of variation, spread over quite a wide dynamic range: rise-times as short as  $\sim 100$  s; flare durations  $\sim 1000$ – $10000$  s; quiescent periods  $\sim 5 \times 10^4$  s long; and a total event duration of more than  $10^7$  s. After  $10^6$  s the best power-law fit to the light curve is that the observed flux decays  $\propto t^{-4/3}$ , but its large variability and spectral changes mean that it might still be consistent with the expected  $t^{-5/3}$  bolometric decay.

Still another surprise came when BAT data *preceding* the trigger were examined. Approximately 3 days earlier, there was a precursor whose peak flux was  $\sim 0.07$  times the

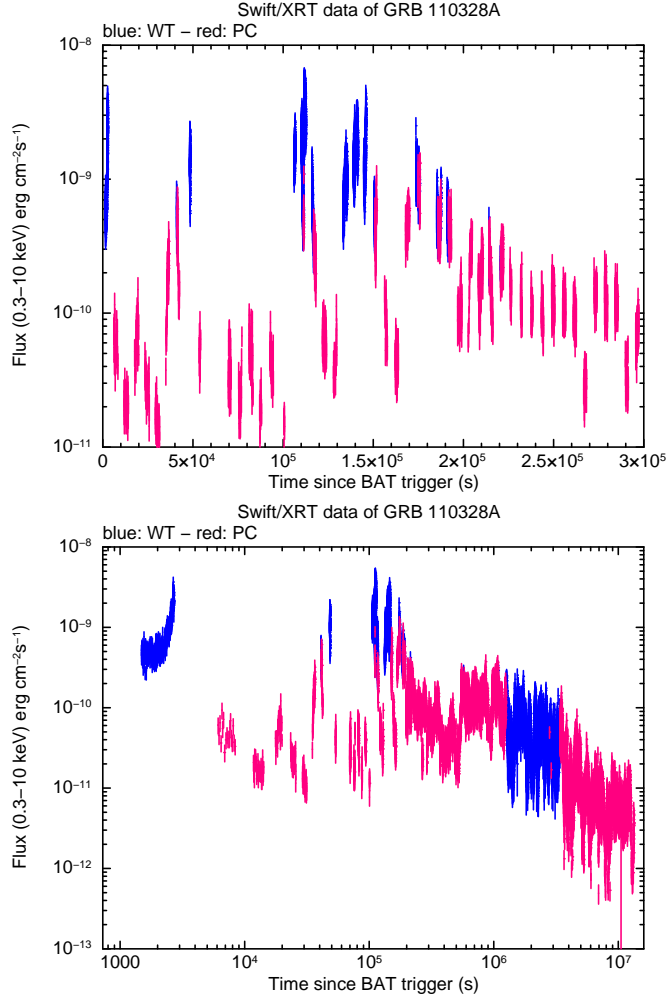


Fig. 1.— Long-term Swift XRT light curves in the 3–10 keV band (blue: WT, red: PC). (Top) Linear in time representation of the first 300,000 s, illustrating the recurring brief flares that gradually widen. (Bottom) Logarithmic in time representation of the entire light curve as of 29 August 2011, five months after activity began. Both these light curves and those in the following figure were generated using the graphing tools of the online Swift Data Repository (Evans *et al.* 2007).

flux of the first flare (Burrows *et al.* 2011). Searches for earlier episodes of emission from the same source in archival data revealed only upper limits.

Optical, near-infrared, and radio observations combine to show that the source lies within 150pc of the center of a galaxy at  $z = 0.354$  (Levan *et al.* 2011). The peak luminosity associated with these flares (interpreted as isotropic) is then  $\simeq 4 \times 10^{48} \text{ erg s}^{-1}$  (Burrows *et al.* 2011). Even at its faintest detected level, the (isotropic) luminosity is  $\simeq 1 \times 10^{46} \text{ erg s}^{-1}$  (Burrows *et al.* 2011). Although relativistic beaming that could enhance the flux by a factor  $\sim 100$  is likely (Burrows *et al.* 2011), the beaming-corrected power is still very large:  $\sim 4 \times 10^{46} \text{ erg/s}$  at the peak, and a total emitted energy  $\sim 5 \times 10^{51} \text{ erg}$  within the first  $1 \times 10^7 \text{ s}$  of the event if the XRT flux is  $\sim 1/3$  of bolometric, as suggested by Bloom *et al.* (2011).

During flaring episodes, the observed  $\nu F_\nu$  is greatest between 10 and 100 keV; it appears to be strongly absorbed by interstellar gas below  $\simeq 5 \text{ keV}$  in the rest-frame. By the time the near-IR measurements were performed, 2–4 days after the event began, their fluxes were  $\sim 10^{-4}$  of the hard X-ray flux at its peak, and  $\sim 10^{-2}$  of its flux during quiescent periods (Burrows *et al.* 2011).

Within six weeks of its discovery, several models were proposed to explain it, most concentrating on a picture in which a main sequence star was tidally disrupted by passing too close to a  $10^6$ – $10^7 M_\odot$  black hole (Bloom *et al.* 2011; Burrows *et al.* 2011; Cannizzo, Troja & Lodato 2011; Socrates 2011). Although these models differ from one another in some of their details, they share a common outline: A main sequence star is thoroughly disrupted as it passes near a black hole, and much of its mass is distributed into an accretion disk around the black hole. A powerful jet is then created, whose radiation, created by the synchro-Compton mechanism, dominates what we observe. In this picture, the rise time of the flares is interpreted as indicating (after allowance for relativistic effects)

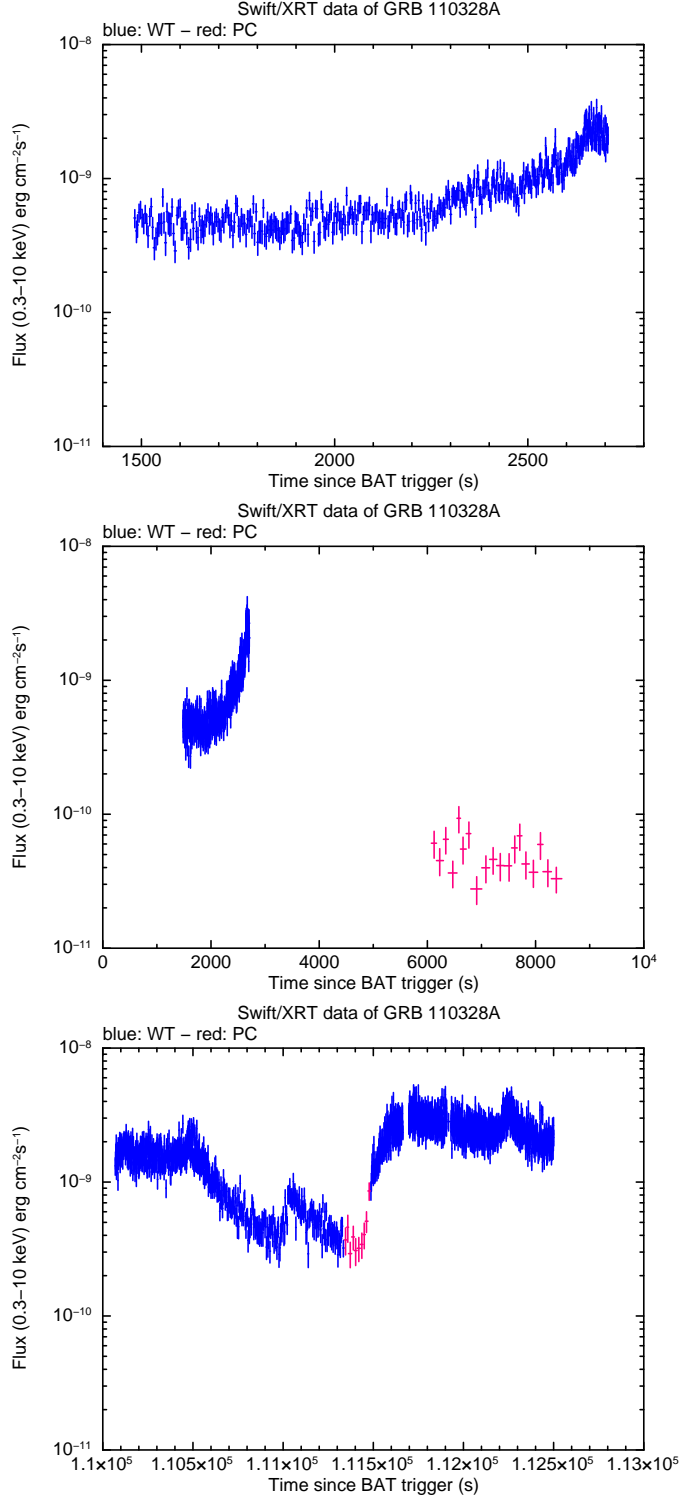


Fig. 2.— Detailed time structure of the flares. (Top) Short timescale structure of the first flare. (Middle) Overview of the first flare. (Bottom) Short timescale structure in the flare  $\simeq 110,000$  s after the BAT trigger. Note that the entire span of the data in the bottom panel is only 1500 s.

the size of individual knots in the jet, and these are related to the gravitational radius of the black hole. The duration of the entire event is thought to indicate either the expected  $t^{-5/3}$  scaling (Burrows *et al.* 2011) due to the orbital period distribution of tidal streams (Rees 1988; Phinney 1989) or the inflow time of the accretion disk (Socrates 2011).

Although plausible in many respects, this consensus model leaves a number of questions unanswered. The tidal disruption radius in units of the black hole gravitational radius is

$$R_T/R_g = R_*(k/f)^{1/6}(M_{BH}/M_*)^{1/3}/(GM_{BH}/c^2) = 50(k/f)^{1/6}M_{BH,6}^{-2/3}\mathcal{M}_*^{2/3}, \quad (1)$$

where  $k$  is the apsidal motion constant (determined by the star’s radial density profile) and  $f$  is its binding energy in units of  $GM_*^2/R_*$  (Phinney 1989). Here  $\mathcal{M}_*$  is the mass of the star in solar units, and we have taken the approximate scaling that  $R_* \propto M_*$  on the main sequence. The ratio  $k/f$  is  $\simeq 0.02$  for radiative stars, but 0.3 for convective stars (Phinney 1989). The circular orbital period at the tidal radius is then proportional to  $\mathcal{M}_*$  and *independent* of  $M_{BH}$  because the dynamical time at the tidal radius is always  $\sim (G\rho_*)^{-1/2}$ , and on the main sequence the mean stellar density  $\rho_* \propto M_*^{-2}$ :

$$P_{\text{orb}}(R_T) = 1.0 \times 10^4 \mathcal{M}_* \text{ s}. \quad (2)$$

However, the *actual* orbital period of matter captured in tidal disruption is likely to be considerably longer. As pointed out by Rees (1988), if the disrupted star approached on a nearly-parabolic orbit, the tidal streams follow highly-eccentric elliptical orbits with a wide range of energies, and therefore of semi-major axes and orbital periods. Their mean energy is likely to be comparable to the star’s self-gravitational binding energy because the gravitational force of the black hole does that much work expanding and disrupting the star. On the other hand, tidally-induced rotation and the gradient of the gravitational potential across the star can lead to some gas being trapped on orbits with semi-major axis as small as  $a_d \sim R_T(M_{BH}/M_*)^{1/3}$ , even while the typical stream’s orbit is larger by a

factor  $(M_{BH}/M_*)^{1/3}$ . It follows that even the shortest orbital period is larger than that of a circular orbit at  $R_T$  by  $\sim (M_{BH}/M_*)^{1/2}$ , here a factor  $\sim 10^3$ , and the typical period likely another factor of  $\sim 10^3$  longer than that. Intersections between stream orbits could lead to conversion of orbital energy to heat, diminishing these orbital periods, but in no case would they become shorter than  $P_{\text{orb}}(R_T)$ . In fact, numerical simulation of the disruption of a main sequence star by a  $10^6 M_\odot$  black hole (Ayal, Livio & Piran 2000) shows a rather continuous accretion rate with a rise time of a few times  $10^5$  s and an overall duration of a few times  $10^6$  s, as expected from these analytic estimates.

It is hard to reconcile these timescales to those seen in the lightcurve. The rise time in the consensus model is said to reflect the light-crossing time across the black hole’s horizon, but it is not clear what dynamics link that quantity to triggering a flare, nor is there any natural explanation for the flare duration. The circular orbital period at  $R_T$  is comparable to the inter-flare interval, but the orbital period of the tidal streams following eccentric orbits, which is the timescale at which the  $t^{-5/3}$  decay begins, appears to be at least two orders of magnitude longer than the time at which the flares merged into a smoother lightcurve. These difficulties have led some (e.g., Cannizzo, Troja & Lodato (2011)) to pose special requirements on this model. They suggest that the pericenter distance must not be a great deal larger than the black hole’s ISCO, so that inflow is largely dynamical. If so, the black hole mass is  $\sim 10^7 M_\odot$ . By the estimate of Equation 1, the pericenter distance must then be a small fraction of  $R_T$ .

Another question is how exactly the accretion drives a jet. If, as is generally believed, jets associated with black holes are powered by some variant of the Blandford-Znajek mechanism (Blandford & Znajek 1977; McKinney 2005; Hawley & Krolik 2006), substantial magnetic field must be attached to the black hole horizon. The Poynting luminosity in the jet is then  $\sim \phi B^2 r_g^2 c$ , where  $\phi$  is a dimensionless quantity that depends on the field geometry



and increases with black hole spin parameter  $a/M$ , but is generally significantly less than unity. In this object, the field on the horizon must then be  $\sim 1 \times 10^6 \phi^{-1/2} L_{45}^{1/2} M_{BH,6}^{-1}$  G. When the system has a long lifespan (e.g., in AGN), the magneto-rotational instability stirs MHD turbulence in the accretion disk and builds the magnetic field; numerical simulations of MRI-driven MHD turbulence show that generically  $\sim 10$  orbital periods are required to reach saturation (Stone *et al.* 1996). Accretion may also accumulate magnetic flux on the horizon (Beckwith, Hawley & Krolik 2009); its build-up rate depends, of course, on the structure of the large-scale magnetic field. Here, the accretion flow must plunge into the black hole on a dynamical time, and it is unclear whether such a strong field could be generated.

Still another problem raised by this model is how to understand the very large amplitude and very rapidly-varying flares. As recognized by Bloom *et al.* (2011), relativistic jets in blazars behave very differently: Their characteristic fluctuation amplitudes are more typically factors of a few than factors of a few orders of magnitude, and their duty cycle at high flux is usually considerably greater than it is during the first several flaring episodes of this event. One could rephrase this question as, “What generates the extremely bright, compact, and short-lived knots that in this model are assumed to account for the flares?”

In this paper, we propose an alternative version of the tidal disruption model that we believe holds some promise for supplying answers to all of these questions. We suggest that the star that is tidally disrupted is a white dwarf, not a main sequence star, and that it is not disrupted all at once, but instead loses pieces of itself in several passes before dissolving (see Fig. 3 for a schematic cartoon) an idea previously considered in the context of white dwarfs on more nearly circular orbits by Sesana *et al.* (2008) and Zalamea, Menou & Beloborodov (2010). This suggestion is motivated by the fact that the fundamental timescale of a tidal disruption is dictated by the mean density of the star;

the greater density of a white dwarf makes it much easier to achieve the short timescales of this event. In both our picture and the main sequence star model, the accretion flow drives a jet, which produces the observed radiation; as we shall see, this, too, is favored quantitatively by the higher density of a white dwarf. The remainder of this paper will develop the consequences of these ideas.

Note that tidal disruption of a white dwarf by a black hole has been also previously discussed in the context of the possible nuclear ignition of the white dwarf (Wilson & Mathews 2004; Dearborn, Wilson & Mathews 2005; Rosswog, Ramirez-Ruiz & Hix 2008). This requires a rather small black hole and a deep encounter, with pericenter a small fraction of  $R_T$ . Here we focus on larger black holes and more distant tidal disruptions that do not lead to such an explosion.

## 2. Timescales

Using the mass-radius relations of Nauenberg (1972), we find that a white dwarf is disrupted at

$$R_T/R_g \simeq 8M_{max}^{1/3}\mathcal{M}_*^{-2/3}M_{BH,4}^{-2/3}\left[1. - 0.64(\mathcal{M}_*/M_{max})^{4/3}\right], \quad (3)$$

where  $M_{max} \simeq 1.4$  is the maximal mass of a white dwarf in units of  $M_\odot$ . In this estimate, we used an apsidal motion constant  $k = 0.14$  (Sirotkin & Kim 2009) and a binding energy factor  $f = 6/7$ , both appropriate to an  $n = 3/2$  polytrope. Note that we have changed our fiducial black hole mass from  $10^6$  to  $10^4 M_\odot$  because the higher densities of white dwarfs require smaller mass black holes in order to keep  $R_T/R_g > 1$ . In fact, white dwarf tidal disruption requires  $M_{BH} < 2 \times 10^5 M_\odot$ , and somewhat less than that if the black hole rotates slowly or it is desired that  $R_T$  exceed the ISCO. The period of a circular orbit at

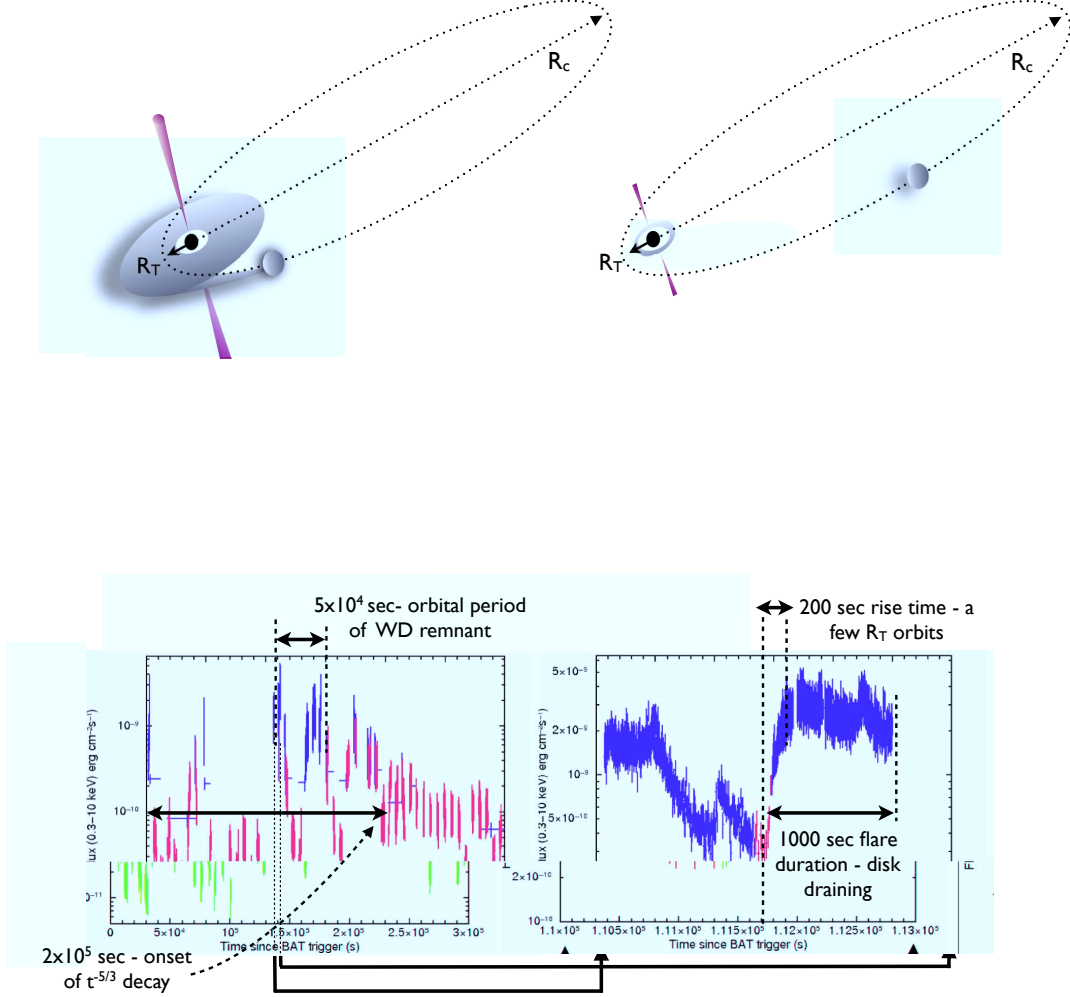


Fig. 3.— Top: A schematic description of our model, a white dwarf in a highly eccentric orbit that passes near a massive black hole. A significant fraction of the white dwarf mass is torn off around the tidal radius,  $R_T$ . The debris produces a small but massive accretion disk that supports a jet. Once the white dwarf moves away from the black hole the tidal disruption ceases, the disk is drained, and the jet dies out. Bottom: Different observed time scales and their relation to the schematic model.

the white dwarf tidal radius is (ignoring black hole spin)

$$P_{\text{orb}}(R_T) \simeq 6\mathcal{M}_*^{-1} \text{ s}, \quad (4)$$

three orders of magnitude shorter than for a disrupted main sequence star. In this regard, it is worth noting that the white dwarf mass distribution appears to be centered at  $\simeq 0.6\text{--}0.7M_\odot$  (Hansen & Liebert 2003).

Matter torn from a white dwarf, like that taken from a main sequence star, will travel initially on a wide range of highly-eccentric orbits. Following the arguments of Rees (1988), the most tightly-bound matter might be expected to have an orbital period  $\sim 900M_{BH,4}^{1/2}\mathcal{M}_*^{-3/2}$  s. However, there are several potential mechanisms that could substantially shorten this period. If the black hole mass is not greatly smaller than its upper bound, the orbital pericenters will be in the strongly relativistic regime. There, as pointed out by Cannizzo, Lee & Goodman (1990), because relativistic effects cause differential pericenter precession, different stream orbits can intersect, so that dissipation transfers energy from orbital motion to heat and possibly photons. This process may be able to circularize the orbits on a timescale of a few orbital periods. When the black hole spins rapidly and its mass is close to the upper bound for tidal disruption,  $R_T/R_g$  may be small enough that the orbital precession rate becomes comparable to the orbital frequency, enhancing this effect (we thank the anonymous referee for this suggestion).

Multi-pass tidal break-ups may also lead directly to tighter orbits for tidally-disrupted matter. The orbital time of the white dwarf remnant is  $\sim 10$  disk drainage times, so the only mass remaining in the disk by the time the remnant returns is the mass pushed outward absorbing the angular momentum of the accreted mass. If the mass remaining in the disk from the previous encounter is a fraction  $\epsilon$  of the newly-arriving mass, the momentum of the matter breaking off the star near  $R_T$  will be reduced by that fraction and its kinetic energy reduced by  $2\epsilon$ . The semi-major axis of the resulting orbit for matter

with the mean incoming orbital energy (i.e., near zero) is then  $R_T/(4\epsilon)$ , giving an orbital period  $P_{\text{orb}}(R_T)(4\epsilon)^{-3/2}$ . If  $\epsilon$  is as little as 0.05, the resulting period is only  $\sim 10P_{\text{orb}}(R_T)$ . Because there was a precursor to this event, it is possible that this sort of process may have occurred as early as the flare that caused the BAT trigger.

As remarked before, internally-generated magnetic field requires  $\sim 10$  orbital periods to reach saturation. When the period is as short as  $P_{\text{orb}}(R_T)$  for a white dwarf tidal disruption, the magnetic field can grow to full strength within  $\sim 100$  s. Such rapid growth would then be consistent with explaining the rapid rise time of the flares.

The  $\sim 1000$ – $10000$  s duration of individual flares might now be reinterpreted as the drainage time for the small, temporary accretion disk formed from captured material. Once the turbulence has reached saturation, this timescale  $t_{\text{in}} \sim (P_{\text{orb}}/2\pi)\alpha^{-1}(R_T/H)^2$  if matter moves inward only by the action of internal disk stresses, where  $\alpha \sim 0.1$  is the ratio of vertically-integrated magnetic stress to vertically-integrated pressure, and the disk density scale height is  $H$ . Close to the ISCO, however, the mean stress overestimates the inflow time because stress fluctuations can remove enough angular momentum from fluid elements to send them all the way into the plunging region, permanently removing them from the disk. As a result, Krolik, Hawley & Hirose (2005) found that for  $R < 3R_{\text{ISCO}}$  in a disk with saturated MHD turbulence,  $t_{\text{in}} \sim 10P_{\text{orb}}(R)$ . Allowing for some material on orbits with periods greater than  $P_{\text{orb}}(R_T)$ ,  $\gtrsim 1000$  s might then be a reasonable estimate for the total duration of a single flare, including both build-up of the magnetic field and inflow.

If the initial pericenter of the white dwarf’s orbit is  $\sim R_T$ , in its first passage it may lose only a fraction of its mass (indeed the detection of a precursor flare  $\sim 3$  days before the BAT trigger [Burrows *et al.* (2011)] suggests that the very first passage may have been slightly outside  $R_T$ ). If that is the case, some of its mass is captured into orbiting streams by the black hole, but the rest of the star remains a coherent bound object, albeit

substantially stretched and distorted. It can then make several more passes through the region near  $R_T$ , traversing a highly eccentric, but well-defined, orbit.

If its initial approach was on a parabolic orbit, the white dwarf’s net binding energy to the black hole is likely to be comparable to its initial self-gravitational binding energy because the energy to alter its structure is taken from the orbital energy. The period of its orbit after capture is then approximately

$$P_{\text{orb}}(R_c) \sim \frac{2\pi}{(2f)^{3/2}} \left( \frac{M_{BH}}{M_*} \right) \left( \frac{R_*^3}{GM_*} \right)^{1/2} \sim 6 \times 10^4 \mathcal{M}_*^{-2} M_{BH,4} \text{ s.} \quad (5)$$

It is natural to identify this period with the interval between flares, as each time the star passes through the pericenter of its orbit it loses a fraction of its mass (see Fig. 3). Moreover, because in each of these passes additional gravitational work is done stretching the remnant of the white dwarf, its orbit changes from pericenter passage to pericenter passage, explaining the irregularity of the inter-flare intervals. These structural changes may also alter the effective  $R_T$  from one passage to the next.

When the white dwarf is finally disrupted completely (presumably just before the final flare at  $\sim 1.7 \times 10^5$  s after the initial BAT trigger), its remaining mass finds itself spread over orbits with a wide range of binding energies and periods. The tidal streams with orbital periods longer than the remnant’s orbital period return to the neighborhood of  $R_T$  only after a comparatively long time, extending the duration of the event. If the distribution function of tidal-stream mass with orbital energy is roughly flat (Rees 1988; Phinney 1989; Lodato, King & Pringle 2009), the late-time accretion rate should decline  $\propto t^{-5/3}$  from the time at which the white dwarf is completely disrupted onward.

### 3. Accretion physics and driving the jet

In the interior of a white dwarf, the electrons are highly degenerate. The equivalent temperature of the Fermi level  $E_F$  is  $2 \times 10^9 \rho_6^{2/3}$  K, where we scale to a characteristic density of  $10^6 \text{ gm cm}^{-3}$  because the mean density of a white dwarf is  $1.1 \times 10^6 \mathcal{M}^2 M_{max}^{-1} / (1 - 0.64(\mathcal{M}/M_{max})^{4/3}) \text{ gm cm}^{-3}$  (again using the mass-radius relation of Nauenberg (1972)). On the other hand, typical interior temperatures are  $\sim 10^7$  K (Hansen 2004). The initial phases of tidal disruption do not change this degree of degeneracy because adiabatic expansion leaves the ratio  $E_F/kT$  invariant (Landau & Lifshitz 1980).

The flow does not remain adiabatic for long, however. For example, orbital precession leads to shocks. In these shocks, electron heating will initially be retarded by scattering-suppression due to their degeneracy, but there is no such constraint on the ions. The slowest ion-electron heating rate occurs when the ions are so much hotter than the electrons that the relative velocity is dominated by the ion thermal speed. Before allowing for electron degeneracy, the characteristic time for heat transfer by ion-electron Coulomb scattering in the disk formed by tidal disruption of a white dwarf is

$$t_{ion,heat} \sim 8 \times 10^{-6} (H/R) M_{BH,4} T_{i,10}^{3/2} \mathcal{M}_*^{-3} (\Delta M/M_*)^{-1} \text{ s.} \quad (6)$$

Here we have set the Coulomb logarithm to 30 and assumed that, appropriate to a C/O composition, the mean ion mass is  $14m_p$ . The ratio  $\Delta M/M_*$  is the fraction of the white dwarf mass deposited in a disk with radius  $R_T$  and scale height  $H$ . Degeneracy retards energy transfer to electrons because the crowded phase space partially suppresses collisions. Only electrons with initial momenta close enough to the Fermi level momentum that the momentum transfer per event  $\Delta p_{ie} \sim m_e \langle v_i \rangle$  can lift them to the region of unoccupied states can participate. The fraction of electrons in the total population able to scatter is then

$$f_{scatter} \sim 3(m_e/2m_i)^{1/2} (kT_i/E_F)^{1/2} \sim 0.05 (kT_i/E_F)^{1/2} \quad (7)$$

if  $\Delta p_{ie} \ll p_F$ . The unadjusted heating time is so short that it is hard to imagine circumstances in which the degeneracy correction could make any difference on the timescales relevant to this situation.

Thus, the electrons and ions can be expected to thermally equilibrate very rapidly, and at these densities and temperatures, the electrons (now no longer degenerate) will also rapidly thermally equilibrate with radiation. However, the radiative cooling time of the system is much longer than any of the relevant timescales:

$$t_{\text{cool}} \sim 1 \times 10^{11} (H/R) (\Delta M/M_*) \mathcal{M}_*^{7/3} M_{BH,4}^{-2/3} \text{ s.} \quad (8)$$

The radiation pressure is therefore effectively trapped within the material. Because shrinking the highly eccentric orbits to nearly circular requires dissipating an energy comparable to the orbital energy, the disk can therefore be expected to be geometrically thick,  $H/R \sim 1$ . This is the regime of photon-trapping associated with super-Eddington accretion Begelman (1979); Abramowicz *et al.* (1988). Because the diffusion time is long compared to the inflow time, the radiation intensity distribution is far from steady-state, and the emergent luminosity is much less than the the rate at which heat is dissipated. Consequently, the ratio of thermal disk luminosity to rest-mass accretion rate is much less than the conventional  $\sim 0.1 \dot{M} c^2$ .

The total pressure in the inner disk can be expected to be of order the electron density times the local virial temperature,

$$p_{\text{disk}} \sim 5 \times 10^{21} (\Delta M/M_*) \mathcal{M}_* M_{BH,4}^{-3} (H/R)^{-1} \left( \frac{R}{10R_g} \right)^{-4} \text{ dyne cm}^{-2}. \quad (9)$$

If the magnetic pressure on the event horizon is limited by the inner disk pressure (the simulations of Beckwith, Hawley & Krolik (2009) suggest that it may be a factor of 3–4 smaller), the expected field strength would be

$$B_{\text{hor}} \lesssim 4 \times 10^{11} (\Delta M/M_*)^{1/2} \mathcal{M}_*^{1/2} M_{BH,4}^{-3/2} (H/R)^{-1/2} \left( \frac{R}{10R_g} \right)^{-2} \text{ G.} \quad (10)$$



One might then predict a jet Poynting power as much as

$$L_{\text{jet}} \sim 3 \times 10^{50} \phi (B_{\text{hor}}^2 / 8\pi p_{\text{disk}}) (\Delta M / M_*) \mathcal{M}_* M_{BH,4}^{-1} \text{ erg/s.} \quad (11)$$

Even after allowing for a field rather less intense than the disk pressure,  $\Delta M / M_* < 1$ , and a small coefficient  $\phi$ , it would seem that the jet power could easily reach the level seen (peak power after allowance for beaming of  $\sim 3 \times 10^{46} \text{ erg s}^{-1}$ ). Note, also, that  $L_{\text{jet}} \propto M_{BH}^{-1}$ , so models requiring larger black holes tend to generate weaker jets. As pointed out by our anonymous referee, the smaller (in  $R_g$  terms) disks characteristic of white dwarf disruptions also minimize Compton drag (Phinney 1987).

Numerical general relativistic MHD simulations of accretion have shown that the luminosity of the jet can also be estimated in terms of an effective “efficiency” per unit rest-mass accreted that is a function of the black hole spin (McKinney 2005; Hawley & Krolik 2006). Using the expression found in the latter reference, one might predict

$$L_{\text{jet}} \sim 4 \times 10^{49} [1 - |a/M|]^{-1} (\Delta M / M_*) \mathcal{M}_*^2 \text{ erg/s} \quad (12)$$

if the inflow time is  $\simeq 10P_{\text{orb}}(R_T)$ . For a black hole with spin parameter  $a/M \simeq 0.9$ , this estimate agrees with the previous one evaluated for the fiducial parameters. That it should do so is no coincidence—as assumed in the previous estimate, the magnetic field intensity on the horizon is comparable to the inner disk pressure in these simulations. However, it should be borne in mind that the energy for this jet is actually drawn from the reducible mass (the rotational kinetic energy) of the black hole, not the accretion flow. The function of the accretion is solely to sustain a strong magnetic field on the black hole horizon.

Although nominally independent of black hole mass, in fact this second estimate has an implied dependence through the bound on black hole mass placed by the density of the disrupted star. In rough terms, the second jet luminosity estimate simply mirrors the accretion rate, which is  $\propto \Delta M \Omega(R_T)$ . In both the main sequence star and white dwarf

models, the amount of mass placed in orbit is  $\sim M_\odot$ . Where they differ is exactly the point we have emphasized in the context of the lightcurve’s timescales:  $\Omega(R_T) \sim (G\rho_*)^{1/2}$ , which is  $\sim 10^3$  times larger for white dwarfs than for main sequence stars.

When the disk is drained, the inner disk pressure falls. If that permits the flux on the horizon to expand, so that the magnetic field there becomes weaker, the jet would be correspondingly diminished. Without continuing accretion—that is, between episodes of tidal capture—radiation from the jet would be much reduced. This is, of course, what one would expect in those intervals when the remnant of the white dwarf is out near apocenter.

#### 4. Summary

We propose that the remarkable event known as Swift J1644+57 is more likely the tidal disruption of a white dwarf than a main sequence star. The fact that white dwarfs are typically  $\sim 10^6$  times denser than main sequence stars makes it much easier to understand the very short timescales characteristic of this object’s lightcurve. Postulating that the initial encounter was not close enough to disrupt the white dwarf completely explains the remarkable flares seen over the event’s first few days (see Fig. 3). The  $\sim 100$  s fluctuations during the flares may then be identified with the inflow timescale from radii  $\sim R_T$ , which in this case is likely only  $\sim 10R_g$ . This timescale is imprinted on the lightcurve if the radiation emerges from the jet when it has traveled less than a few hundred seconds (in the observer’s frame) from the black hole, a distance equivalent to  $\sim 10^3$ – $10^4 R_g$ . The disk drainage time is longer because the disk can be expected to spread to radii somewhat larger than  $R_T$ ; this accounts for the flare durations,  $\sim 10^3$ – $10^4$  s. The interflare time  $\sim 5 \times 10^4$  is the orbital period of the white dwarf remnant. Because the remnant orbital period is rather longer than the drainage time, there is an extended period after most of the disk mass has been accreted, but before it is refilled by the next pericenter passage of the white dwarf, when

the accretion rate is low, and the jet is therefore weak.

The higher densities of white dwarfs also explain both why the observed radiation is dominated by a jet and the jet’s high luminosity. The very high surface density of the disk thoroughly traps thermal photons, while the associated high pressure can support a stronger magnetic field on the black hole. Although the jet luminosity associated with a main sequence star event is nominally just enough to supply the observed radiated power (after beaming corrections), that may still be inadequate. The numbers quoted refer only to luminosity in observed bands; it is possible there is additional luminosity elsewhere in the spectrum (e.g., between 100 keV and 100 MeV). More importantly, in most jet radiation models (e.g., Celotti & Ghisellini (2008)), the photon luminosity is only  $\sim 1\% - 10\%$  of the kinetic power; the jet powers estimated on the basis of the disk pressure or accretion rate refer to the total, so that  $L_{\text{jet}}$  must be significantly *greater* than the observed luminosity.

Although the space density of white dwarfs ( $\simeq 3 \times 10^{-3} \text{ pc}^{-3}$  in the Solar neighborhood: Rowell & Hambly (2011)) is very similar to the space density of solar-mass main sequence stars ( $\simeq 3.5 \times 10^{-3} \text{ pc}^{-3}$ : Reid, Gizis & Hawley (2002)), their effective cross section for coming close enough to a black hole to be tidally disrupted is smaller by the ratio of their tidal disruption radii,  $\simeq 2 \times 10^{-3} (\mathcal{M}_{*,\text{wd}}/\mathcal{M}_{*,\text{ms}})^{-2/3}$ . Burrows *et al.* (2011) estimated the total number of solar-mass main sequence star disruptions within the volume detectable by Swift to be  $\sim 10^4 \text{ yr}^{-1}$ , suggesting that the rate of white dwarf tidal disruptions should be  $\sim 10 \text{ yr}^{-1}$ . If they are all relativistically beamed so that we see only  $\sim 1\%$  of all events, the observable rate falls to  $\sim 0.1 \text{ yr}^{-1}$ , and is reduced further by a factor  $\sim 10$  to account for BAT’s field of view and observing efficiency (Burrows *et al.* 2011). Thus, we might expect BAT to detect perhaps  $\sim 0.01 \text{ yr}^{-1}$ , rather than the actual  $\sim 1/6 \text{ yr}^{-1}$ . On the other hand, the rates estimated by Burrows *et al.* (2011) (and similarly by Zauderer *et al.* (2011)) were based on theoretical predictions that are an order of magnitude smaller than

some empirically-based rate estimates (Gezari *et al.* 2008; Maksym, Ulmer & Eracleous 2010). Given the many uncertainties (and small-number statistics), we believe the rate of white dwarf disruptions is consistent with the Swift detection rate. Moreover, a large part of the contrast between our predicted rate and the rate expected for main sequence stars would be removed if the required pericenter distance for a main sequence tidal disruption is  $\sim 0.1R_T$ , as suggested by Cannizzo, Troja & Lodato (2011).

We close with a final significant contrast between white dwarf and main sequence star models for this event. The black holes most effective at tidal disruption of white dwarfs are smaller by the ratio  $(\rho_{*,\text{ms}}/\rho_{*,\text{wd}})^{1/3} \sim 10^{-2}$ . Although there has long been good evidence for black holes of  $\sim 10^6 M_\odot$  or more in the central regions of Galaxies (beginning with the Milky Way) and somewhat shakier evidence for black holes in galactic nuclei with masses  $\gtrsim 10^5 M_\odot$  (collected in Xiao *et al.* (2011), these masses are all based on the assumption that gravitational dynamics dominate broad-line gas motions, and in most cases use only scaling arguments to estimate the broad-line region’s size), this event represents the first indication of a black hole in the  $\sim 10^4$ – $10^5 M_\odot$  mass range. On the basis of the  $M_{BH}$ –bulge luminosity correlation, Burrows *et al.* (2011) estimated that this galaxy should have a nuclear black hole with mass  $\simeq 2 \times 10^7 M_\odot$ . If we are correct in our suggestion that this event was due to disruption of a white dwarf rather than a main sequence star, the correlation estimate may be too large by a factor of 100 or more in this galaxy. Alternatively, the black hole responsible for this disruption could be a second, smaller black hole orbiting a larger one, whose mass is closer to the correlation estimate.

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## REFERENCES

- Abramowicz, M. A., Czerny, B., Lasota, J. P., and Szuszkiewicz, E. 1988, ApJ, 332, 646.
- Ayal, S., Livio, M., and Piran, T. 2000, ApJ, 545, 772, arXiv:astro-ph/0002499.
- Beckwith, K., Hawley, J. F., and Krolik, J. H. 2009, ApJ, 707, 428, 0906.2784.
- Begelman, M. C. 1979, MNRAS, 187, 237.
- Blandford, R. D. and Znajek, R. L. 1977, MNRAS, 179, 433.
- Bloom, J. S. *et al.* 2011, Science, 333, 203, 1104.3257.
- Burrows, D. N. *et al.* 2011, Nature, 476, 421.
- Cannizzo, J. K., Lee, H. M., and Goodman, J. 1990, ApJ, 351, 38.
- Cannizzo, J. K., Troja, E., and Lodato, G. 2011, ArXiv e-prints, 1105.2816.
- Celotti, A. and Ghisellini, G. 2008, MNRAS, 385, 283, 0711.4112.
- Dearborn, D. S. P., Wilson, J. R., and Mathews, G. J. 2005, ApJ, 630, 309.
- Evans, P. A. *et al.* 2007, A&A, 469, 379, 0704.0128.
- Gezari, S. *et al.* 2008, ApJ, 676, 944, 0712.4149.
- Hansen, B. 2004, Phys. Rep., 399, 1.
- Hansen, B. M. S. and Liebert, J. 2003, ARA&A, 41, 465.
- Hawley, J. F. and Krolik, J. H. 2006, ApJ, 641, 103, arXiv:astro-ph/0512227.
- Krolik, J. H., Hawley, J. F., and Hirose, S. 2005, ApJ, 622, 1008, arXiv:astro-ph/0409231.
- Landau, L. D. and Lifshitz, E. M. 1980, Statistical physics, : Oxford: Pergamon Press).

- Levan, A. J. *et al.* 2011, *Science*, 333, 199, 1104.3356.
- Lodato, G., King, A. R., and Pringle, J. E. 2009, *MNRAS*, 392, 332, 0810.1288.
- Maksym, W. P., Ulmer, M. P., and Eracleous, M. 2010, *ApJ*, 722, 1035, 1008.4140.
- McKinney, J. C. 2005, *ApJ*, 630, L5, arXiv:astro-ph/0506367.
- Nauenberg, M. 1972, *ApJ*, 175, 417.
- Phinney, E. S. 1987, in *Superluminal Radio Sources*, ed. J. A. Zensus & T. J. Pearson, 301.
- Phinney, E. S. 1989, in *The Center of the Galaxy*, ed. M. Morris, volume 136 of *IAU Symposium*, 543.
- Rees, M. J. 1988, *Nature*, 333, 523.
- Reid, I. N., Gizis, J. E., and Hawley, S. L. 2002, *AJ*, 124, 2721.
- Rosswog, S., Ramirez-Ruiz, E., and Hix, W. R. 2008, *ApJ*, 679, 1385, 0712.2513.
- Rowell, N. and Hambly, N. 2011, *ArXiv e-prints*, 1102.3193.
- Sesana, A., Vecchio, A., Eracleous, M., and Sigurdsson, S. 2008, *MNRAS*, 391, 718, 0806.0624.
- Sirotkin, F. V. and Kim, W.-T. 2009, *ApJ*, 698, 715, 0904.2939.
- Socrates, A. 2011, *ArXiv e-prints*, 1105.2557.
- Stone, J. M., Hawley, J. F., Gammie, C. F., and Balbus, S. A. 1996, *ApJ*, 463, 656.
- Wilson, J. R. and Mathews, G. J. 2004, *ApJ*, 610, 368, arXiv:astro-ph/0307337.
- Xiao, T., Barth, A. J., Greene, J. E., Ho, L. C., Bentz, M. C., Ludwig, R. R., and Jiang, Y. 2011, *ArXiv e-prints*, 1106.6232.

Zalamea, I., Menou, K., and Beloborodov, A. M. 2010, MNRAS, 409, L25, 1005.3987.

Zauderer, B. A. *et al.* 2011, ArXiv e-prints, 1106.3568.